

Azimuthal asymmetry of J/ψ suppression in non-central heavy-ion collisions

Xin-Nian Wang^a and Feng Yuan^b

^a *Nuclear Science Division, MS 70-319, Lawrence Berkley National Laboratory, Berkeley, California 94720*

^b *Institut für Theoretische Physik der Universität, Philosophenweg 19, D-69120 Heidelberg, Germany*

Abstract

The azimuthal asymmetry of J/ψ suppression in non-central heavy-ion collisions is studied within a dynamic model of J/ψ suppression in a deconfined partonic medium. Within this model, J/ψ suppression in heavy-ion collisions is caused mainly by the initial state nuclear absorption and dissociation via gluon- J/ψ scattering in deconfined partonic medium. Only the second mechanism gives arise to azimuthal asymmetry of the final J/ψ production. We demonstrate that if there is an onset of suppression by quark-gluon plasma (QGP) in the NA50 data, it must be accompanied by the non-vanishing azimuthal asymmetry. Using the same critical density above which the QGP effect enters, we predict the azimuthal asymmetric coefficient v_2 as well as the survival probability for J/ψ at the RHIC energy.

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In the search for quark-gluon plasma (QGP), J/ψ suppression has been proposed as one of the promising signals [1] of the deconfinement in high-energy heavy-ion collisions. Because of the color screening effect in a quark-gluon plasma, the linear confining potential in vacuum that binds two heavy quarks to form a quarkonium disappears so that it can be easily broken up causing suppression of the J/ψ production. The problem in heavy-ion collisions is however complicated by other competing mechanisms such as initial nuclear absorption [2] and hadronic dissociation [3]. While recent precision data from the NA50 [4] experiment at the CERN SPS energies clearly show anomalous suppression unexplained by the normal initial nuclear absorption, there are still much debates about the exact nature of the anomalous suppression [5,6], whether it is caused by the formation of QGP or dissociation by ordinary hadronic matter.

We propose in this letter the study of azimuthal asymmetry of J/ψ production [7] as additional measurements to distinguish different competing mechanism of J/ψ suppression. Since the initial state interactions such as nuclear absorption or nuclear shadowing of gluon distribution has no preference over the azimuthal direction they will not have any contribution to the azimuthal anisotropy of the J/ψ production. Only suppression by the final state interaction with the produced medium will cause significant azimuthal anisotropy in the final J/ψ distribution in the transverse direction. If the centrality dependence of the J/ψ suppression additional to the initial nuclear absorption is caused by formation of QGP, it must be accompanied by a sudden onset of the azimuthal anisotropy. On the other hand, a hadronic absorption scenario would give a continuous centrality dependence of the azimuthal anisotropy. In this letter, we will study the centrality dependence of both the averaged J/ψ suppression factor and the azimuthal anisotropy with a model in which J/ψ suppression is caused by initial nuclear absorption and final state dissociation by QGP above a critical density. Using parameters from fitting the NA50 data, we will also give predictions for J/ψ suppression and its azimuthal anisotropy at the RHIC energies.

We will quantify the azimuthal anisotropy by the second Fourier coefficient v_2 of the azimuthal angle distribution of the final J/ψ distribution, similarly to the proposed elliptic

flow measurement [8,9]. We follow the microdynamic approach of J/ψ suppression [10], in which the J/ψ suppression is caused by gluonic dissociation. Different from a normal hadron gas, a deconfined partonic system contains much harder gluons which can easily break up a J/ψ . The perturbative calculations predict the gluon- J/ψ dissociation cross section [11,12],

$$\sigma_\psi(q^0) = N_0 \frac{(q^0/\epsilon_0 - 1)^{3/2}}{(q^0/\epsilon_0)^5}, \quad (1)$$

where

$$N_0 = \frac{2\pi}{3} \left(\frac{32}{3}\right)^2 \left(\frac{16\pi}{3g_s^2}\right) \frac{1}{m_Q^2}.$$

Here g_s is the strong coupling constant, m_Q is charm quark mass, and q^0 is the gluon energy in the J/ψ rest frame. To break up a J/ψ , q^0 must be larger than the binding energy ϵ_0 . Using similar approach, we have calculated the dissociation cross section for P -wave states by gluons,

$$\sigma_\chi(q^0) = 4N_0 \frac{(q^0/\epsilon_\chi - 1)^{1/2} (9(q^0/\epsilon_\chi)^2 - 20(q^0/\epsilon_\chi) + 12)}{(q^0/\epsilon_\chi)^7}, \quad (2)$$

where ϵ_χ is the binding energy of the P -wave state. For χ_c , it is about 0.250GeV . Because of the small value of the binding energy, the validity of the perturbative calculations for the gluonic dissociation cross section of χ_c might be questionable. So the above formula can only be considered more phenomenological. However, keeping this point in mind, we can see that the above expression still qualitatively reflects the fact that χ_c states are easier to be broken up than J/ψ , because of the much lower energy threshold and the overall larger factor of 4. Therefore, this perturbative calculation gives us a reasonable estimate and guides us to include its contribution for a more complete study of J/ψ suppression in heavy ion collisions. In the following we will consider χ_c contributing to about 40% of the initial J/ψ production and use the above formula to estimate its suppression in a deconfined partonic system.

In the rest frame of a deconfined parton gas, the momentum distribution of thermal gluons will depend on the effective temperature T with an approximate Bose-Einstein distribution, $f(k^0; T) \propto [\exp(k^0/T) - 1]^{-1}$. The velocity averaged dissociation cross sections for charmonia is defined as,

$$\langle v_{rel}\sigma \rangle(T, \vec{p}) = \frac{\int d^3k v_{rel}\sigma(q^0)f(k^0; T)}{\int d^3k f(k^0; T)}, \quad (3)$$

where v_{rel} is the relative velocity between J/ψ and a gluon. In this paper we are only interested in J/ψ production in the midrapidity region, so these cross sections will depend on charmonia transverse momentum p_T as well as the effective temperature T .

With the velocity averaged dissociation cross sections, the survival probabilities of charmonia in the deconfined quark-gluon plasma will have the following form,

$$S^{deconf.}(\vec{b}, \vec{r}, \vec{v}) = \exp\left\{-\int_{\tau_0}^{\tau_f} \langle v_{rel}\sigma \rangle \rho(\vec{r} + \vec{v}\tau, \tau) \Theta(\rho - \rho_c) d\tau\right\}, \quad (4)$$

where τ_0 is the formation time of the quark-gluon plasma, which will be set as $\tau_0 = 1fm$ in the following calculations. The upper limit of the time integral τ_f is determined by the Θ function. Here we introduce the critical density ρ_c above which the QGP dissociation effects enters. We assume it is the same for both J/ψ and χ_c . Because of different binding energies, the effective cross sections will have different temperature dependences for J/ψ and χ_c even if ρ_c is reached. ρ is the local density depending on τ , the initial production point \vec{r} and the velocity \vec{v} of the charmonium particles. The velocity \vec{v} depends on the transverse momentum \vec{p}_T and the azimuthal angle ϕ . The dissociation cross sections $\langle v_{rel}\sigma \rangle$ depend on the effective temperature, which will also depend on the local density $\rho_\tau(b, s)$. In the case of 1+1D Bjorken longitudinal expansion with the initial plasma density $\rho_0 = \rho(\tau = \tau_0)$,

$$\rho_\tau(b, s) = \rho_0(b, s) \left(\frac{\tau_0}{\tau}\right)^\alpha, \quad (5)$$

where $\alpha = 1$. Correspondingly, for the effective temperature,

$$T_\tau(b, s) = T_0(b, s) \left(\frac{\tau_0}{\tau}\right)^{1/3}. \quad (6)$$

For simplification, we relate the local effective temperature with the local density in the following way,

$$T_\tau(b, s) = \kappa (\rho_\tau(b, s))^{1/3}, \quad (7)$$

where $\kappa = [\pi^2/16\zeta(3)]^{1/3}$.

The initial plasma density is related to the rapidity density of gluons,

$$\rho_0(b, s) = \frac{1}{\tau_0} \frac{dN^g}{dyd^2s}(b, s), \quad (8)$$

where d^2s is the transverse area of the overlapping region of two colliding nuclei. We will follow the two-component model [13] and include both the soft and hard contribution to the final hadron production. Assuming that the initial gluon density is proportional to the final hadron rapidity density, we have phenomenologically,

$$\frac{dN^g}{dyd^2s}(b, s) = c[f_b n_b(b, s) + f_p n_p(b, s)], \quad (9)$$

where c is a constant and we will set $c = 1$ in the following calculations, and

$$n_p(b, s) = T_A(s)[1 - \exp(-T_B(b - s)\sigma_{pp})] + T_B(b - s)[1 - \exp(-T_A(s)\sigma_{pp})], \quad (10)$$

$$n_b(b, s) = T_A(s)T_B(b - s)\sigma_{pp}. \quad (11)$$

Since mini-jet cross section at the SPS energy is very small, we will effectively only have the soft contribution which is proportional to the number of participant nucleons. At collider energies such as RHIC the contribution from the hard processes is more important. We will use $f_b = 0.34$ and $f_p = 0.88$ as determined by the PHENIX [14] experiment for $Au + Au$ collisions at $\sqrt{s} = 130$ GeV. We extrapolate these parameters to $\sqrt{s} = 200$ GeV by just multiplying a factor of 1.14 found by PHOBOS [15].

The final expression for the survival probability due to QGP suppression is,

$$S^{deconf.}(\vec{b}, \vec{r}, \vec{v}) = \exp\left\{-\int_{\tau_0}^{\tau_f} \frac{d\tau}{\tau} \langle v_{rel} \sigma_{\psi g} \rangle(\vec{r} + \vec{v}\tau, \tau) \rho_0(b, \vec{r} + \vec{v}\tau) \Theta(\rho - \rho_c)\right\}. \quad (12)$$

It is interesting to note that with a very large constant dissociation cross section the above formula will be equivalent to the model [5] by Blaizot *et al.*, where they assume that all of J/ψ will be dissociated above some critical density. The detailed comparison of this approach and ours will be presented elsewhere.

We can also include the transverse expansion effects on the local parton density following Ref. [16],

$$\rho(\vec{r}, \tau) = \frac{\tau_0}{\tau} \int \frac{d\Omega_{v_t}}{2\pi} \rho_0(\vec{r} - \vec{v}_t \tau), \quad (13)$$

where v_t is the average velocity of the transverse expansion of the parton system. We will use $v_t = 0.4c$ for SPS and $v_t = 0.6c$ for RHIC in the following numerical calculations.

Apart from the above discussed QGP suppression for charmonia states, there is also suppression associated with the initial state interaction, i.e., the nuclear absorption of so-called preresonance of $c\bar{c}$ pairs,

$$S^{abs}(\vec{b}, \vec{r}) = \frac{(1 - \exp(-\sigma_{abs} T_A(r)))(1 - \exp(-\sigma_{abs} T_B(|\vec{b} - \vec{r}|))}{\sigma_{abs}^2 T_A(r) T_B(|\vec{b} - \vec{r}|)}. \quad (14)$$

where σ_{abs} is the absorption cross section of the preresonance with nucleons, for which we will set $\sigma_{abs} = 5.8mb$ in this paper.

By summing up these two contributions, we get the final J/ψ survival probability as

$$S^{sur.}(\vec{b}, \vec{p}_T) = \frac{\int d^2\vec{r} T_A(r) T_B(|\vec{b} - \vec{r}|) S^{abs}(\vec{b}, \vec{r}) S^{deconf.}(\vec{r}, \vec{v})}{\int d^2\vec{r} T_A(r) T_B(|\vec{b} - \vec{r}|)}. \quad (15)$$

From the above expression, we see that the nuclear absorption has no dependence on the azimuthal angle. This means that there is no contribution to v_2 from the initial nuclear absorption. On the other hand, the final state interaction or the dissociation by the deconfined parton gas indeed has azimuthal angular dependence because the parton density is azimuthally asymmetric. Therefore, any finite value of v_2 for J/ψ production should come from the final state interaction. It will provide us important information about the early stage of the quark-gluon plasma.

The azimuthal asymmetric coefficient v_2 then can be calculated as

$$v_2(b, p_T) = \frac{\int d\phi S^{sur.}(\vec{b}, \vec{p}_T) \cos(2\phi)}{\int d\phi S^{sur.}(\vec{b}, \vec{p}_T)}, \quad (16)$$

where ϕ is the azimuthal angle between J/ψ transverse momentum \vec{p}_T and the impact parameter \vec{b} . With this formula we can study both the p_T dependence and the centrality (b) dependence of v_2 .

In the following we present the numerical results of the above approach. We first determine the critical density ρ_c by fitting the SPS data on J/ψ suppression, and then predict the suppression and v_2 at RHIC.

In Fig. 1, we show our results at the SPS energy. The upper plot is the survival probability as a function of transverse energy E_T compared with the experimental data from NA50 [4]. From the fit we determined the critical density $\rho_c = 3.3 fm^{-3}$. The correlation between the impact parameter b and transverse energy E_T [17],

$$P(E_T, b) = \frac{1}{\sqrt{2\pi q^2 a N_p(b)}} \exp \left\{ -\frac{[E_T - q N_p(b)]^2}{2q^2 a N_p(b)} \right\},$$

has been used in the calculations, where we set the parameters as $q = 0.274 GeV$ and $a = 1.27$ [5]. In the calculations, we also include the E_T fluctuation effects as in [5,6,18], which is important for the last few E_T bins. From this figure, we can see that our approach can well describe the experimental data of NA50, and taking into account the transverse expansion improves the fit especially in the peripheral region. As expected, v_2 vanishes at very peripheral collisions, and becomes sizable when the anomalous suppression makes sense. For more central collisions, because of the symmetric geometry of the collisions, the azimuthal asymmetry v_2 vanishes again. Since we assumed that only final state interactions produce finite value of v_2 , its increase with E_T is quite abrupt around the value when the critical density ρ_c is reached. After taking into account the transverse expansion, v_2 is a little higher than the case without transverse expansion. This is quite different from v_2 for jet quenching [16], where transverse expansion is found to reduce v_2 significantly.

The results at RHIC are shown in Fig. 2. The survival probability and v_2 of J/ψ are shown as functions of number of participants. Comparing with the results in Fig. 1, we find that the anomalous suppression enters already at peripheral collisions at RHIC, and the gap between the full suppression and the nuclear absorption alone is much larger than that at SPS. And v_2 is also much larger.

It is interesting to note the different dependences of v_2 on the the transverse expansion at the SPS and RHIC energies. At SPS, it enhances v_2 a little bit while at RHIC it reduces. This is because at these two different energies the dominant suppression sources are different in the region where v_2 is sizable. At SPS, a large part of J/ψ suppression comes from χ_c suppression, because the density is not so high and it is difficult to dissociate directly produced

J/ψ . For this part, the dissociation cross section is very large, in the order of a few mb . So χ_c would be dissociated almost totally when the local density is above the critical value ρ_c , which is more similar to the Blaizot-like model. In this case, the transverse expansion will increase v_2 a little bit because it increases the time duration when χ_c is being dissociated. This increase of time duration even overcomes the decrease of geometrical asymmetry due to transverse expansion leading to an increased v_2 . The dominant suppression, however, comes from directly produced J/ψ (about 60%) at the RHIC energies because of the much higher initial density. The transverse expansion accelerates the decrease of the initial density and reduces the initial geometric asymmetry. This leads to the decrease of v_2 , very similar to the case of jet quenching in dense. Since the final total J/ψ suppression is a mixture of direct suppression and suppression through χ_c , this reduction is not as large as it is for jet quenching [16].

In conclusion, we have studied the azimuthal asymmetry of J/ψ suppression at both SPS and RHIC energies within a dynamic model of charmonia dissociation in a deconfined partonic system. With a critical density $\rho_c = 3.3 fm^{-3}$ we can reproduce well the anomalous suppression found by the NA50 experiment at the SPS. We predicted the azimuthal anisotropy v_2 of the J/ψ suppression at SPS. The existence of a critical density ρ_c for J/ψ suppression leads to a sharp increase of v_2 with E_T , assuming that v_2 only comes from J/ψ dissociation in the deconfined matter. At the RHIC energies, we found that the anomalous suppression already plays a role in peripheral collisions because of high density. In noncentral collisions there is sizable v_2 for J/ψ at $p_T = 3 GeV$.

The experimental study of v_2 for J/ψ suppression will be complimentary to other studies of J/ψ suppression. Together with measurements, such as high p_T hadrons where jet quenching plays an important role [19], these studies will provide valuable information about the early stage of high-energy heavy-ion collisions.

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FIGURES

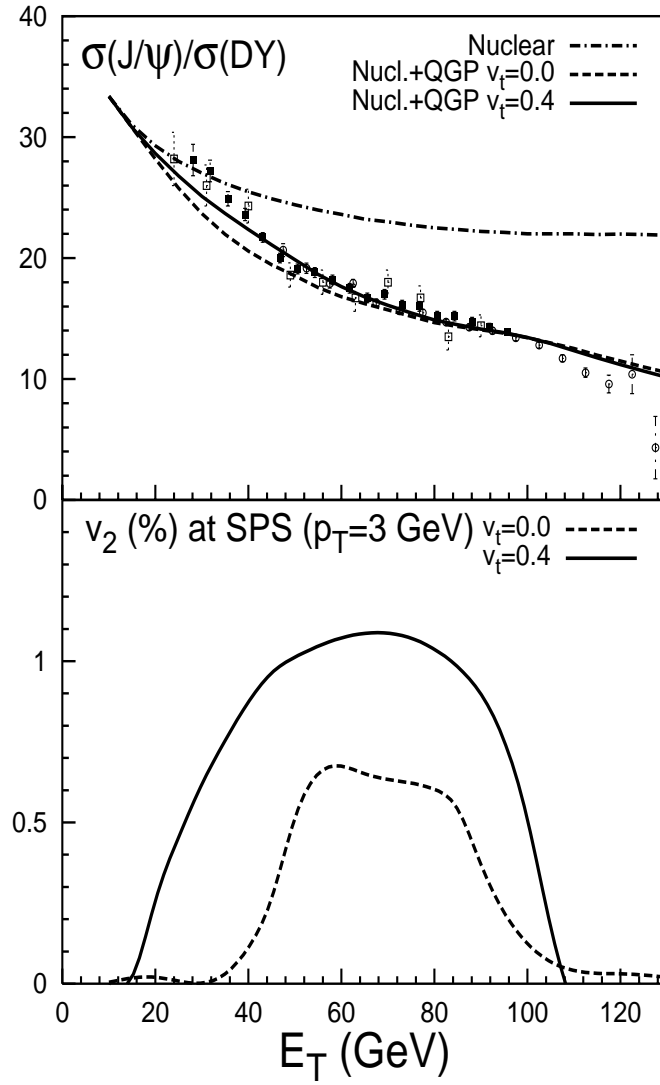


FIG. 1. The J/ψ survival probability and v_2 at SPS as a function of transverse energy E_T : nuclear absorption alone (dot-dashed line); nuclear absorption plus QGP dissociation without (dashed line) and with transverse expansion effects. The experimental data are from NA50 [4].

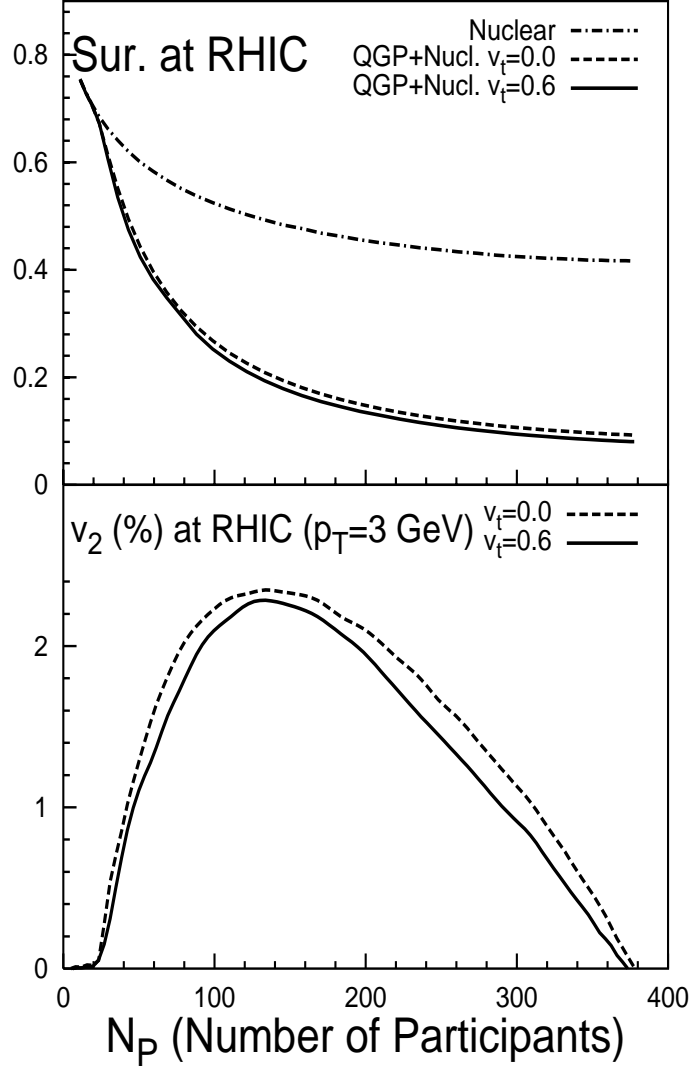


FIG. 2. The J/ψ survival probability and v_2 at RHIC as a function of number of participants N_P .